High-spin 1p-1h configurations in sn-116 and their fragmentation as seen in the reactions sn-116(-]p,p'), sn-116(e,e'), in-115(3he,d) and in-115(alpha, t)

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Published in:
Physics Letters B

DOI:
10.1016/0370-2693(86)91582-0

HIGH-SPIN 1p–1h CONFIGURATIONS IN ^{116}Sn AND THEIR FRAGMENTATION AS SEEN IN THE REACTIONS ^{116}Sn($\vec{p}$, p'), ^{116}Sn(e, e'), ^{115}In(^{3}He, d) AND ^{115}In(α, t)

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Received 16 September 1985; revised manuscript received 14 November 1985

Stretched spin configurations in ^{116}Sn have been studied via the reactions ^{116}Sn($\vec{p}$, p'), ^{116}Sn(e, e'), ^{115}In(^{3}He, d) and ^{115}In(α, t). The high-spin negative parity two-neutron quasiparticle states within the $N = 51-82$ major shell appear to be little fragmented. The most prominent examples are the $I^{\pi} = 9^{-}$ state at $E_x = 3.522$ MeV and two 7^{-} states at 2.909 and 3.120 MeV. It is found that in contrast the proton configurations ($g_{9/2}^{-1}$, $h_{11/2}$) and ($g_{9/2}^{-1}$, $g_{7/2}$) are strongly fragmented. Large-basis BCS shell model calculation, using a SkE-force have been made and DWBA analyses of the (e, e') data and of the ($\vec{p}$, p') data are presented.

It has been of great help in establishing the identity and description of stretched and nearly stretched states to have the same state populated in different reactions. For instance the ($d_{5/2}^{-1}$, $f_{7/2}$)$I^{\pi} = 5^{-}$ and 6^{-} states in ^{28}Si have been studied via inelastic scattering [1] and via proton-stripping reactions [2,3] on ^{27}Al, which is mainly $d_{5/2}^{-1}$ in its ground state. Unfortunately, it is not possible to use this cross-bombardment technique to study stretched states in the heaviest nuclei because the one-hole states of highest j never appear as ground states of possible targets. The most massive target for which such comparison is possible is ^{116}Sn, where the proton $g_{9/2}^{-1}$ configuration is the ground state of ^{115}In.

In order to search for (nearly) stretched states in ^{116}Sn and to study their fragmentation we performed the following experiments:

(i) ^{116}Sn($\vec{p}$, p')^{116}Sn at $E_p = 133.8$ MeV at IUCF, using the QDDM magnetic spectrograph.

(ii) ^{116}Sn(e, e')^{116}Sn at NIKHEF-K, using the QDD spectrometer. Data were taken both at forward and at backward angles, in order to separate longitudinal and transverse form factors.

(iii) ^{115}In(^{3}He, d)^{116}Sn at $E_{^3\mathrm{He}} = 50$ MeV at the

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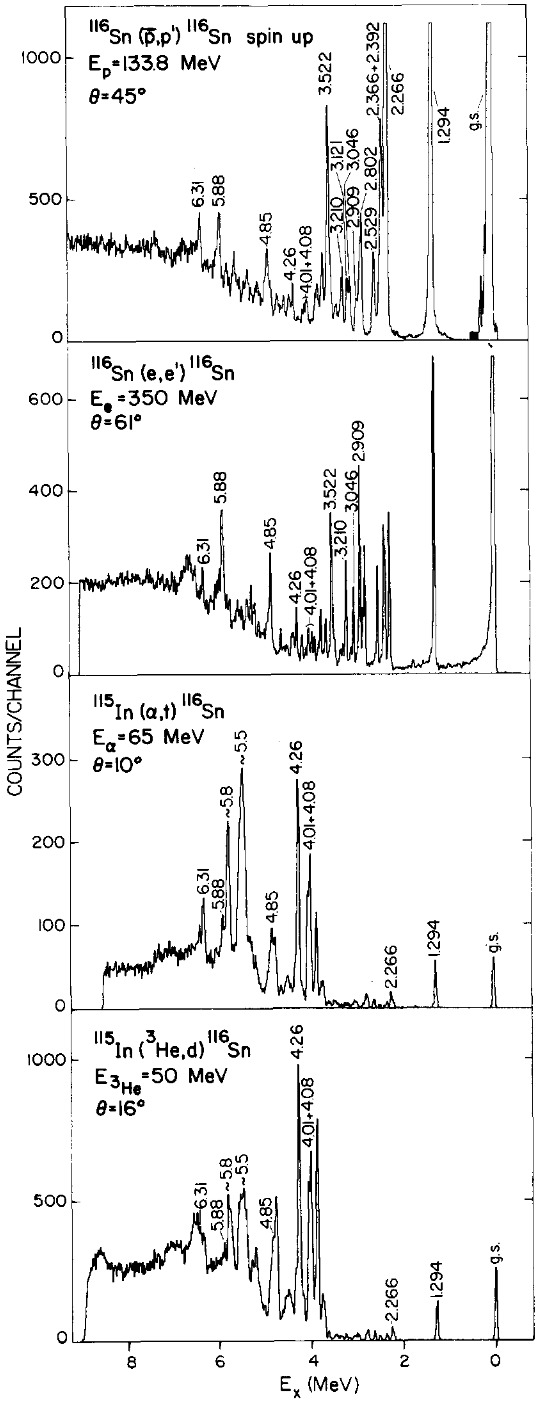

Fig. 1. Experimental spectra of ^{116}Sn observed via, from top to bottom, the ($\vec{p}$, p'), (e, e'), (α, t) and (^{3}He, d) reactions. The excitation energy scale is common for all four spectra.

KVI, using the QMG/2 magnetic spectrograph.

(iv) ^{115}In(α, t)^{116}Sn at E_α = 65 MeV at the KVI, using the same spectrograph.

We present in this letter a first account of the data. Representative spectra from the four different experiments are shown in fig. 1. For both the (p, p') and (e, e') spectra, shown in this figure, the momentum transfer is about 1.9 fm^{-1}, close to the maximum of the form factors of the highest spin states seen. For the (^{3}He, d) a spectrum is shown at an angle that best emphasizes stripping into high-l orbitals. The (α, t) spectrum favours high l-values even more.

The absence of strength below E_x = 3.8 MeV in the stripping spectra illustrates the gap between the Z = 21–50 and the Z = 51–82 major shells. The only appreciable strength is on the gs and the 2$^+$ state at 1.294 MeV and corresponds mostly to the filling of the last hole in the $g_{9/2}$ shell. These features have been noted in earlier work on the ^{115}In(^{3}He, d)^{116}Sn reaction, by Conjeaud et al. [4], Biggerstaff et al. [5] and by Shoup et al. [6]. Those experiments, performed at lower incident energies mainly excite l = 0 and 2 transitions. In the present study we are interested in the $g_{7/2}$– and $h_{11/2}$– shells, which lead to proton particle–hole configurations with maximum spins 8$^+$ and 10$^-$. At our incident energy of 50 MeV the l-matching for these orbitals is such that one may expect to locate the full strengths associated with these orbitals.

The highest spin states observed below the gap in the ($\vec{p}$, p') and (e, e') reactions, are the two I^π = 7$^-$ states at E_x = 2.909 and 3.210 MeV and the I^π = 9$^-$ state at E_x = 3.522 MeV. All three states are known from in-beam gamma work [7].

In the simplest approach the 9$^-$ state is interpreted as a ($g_{7/2}$, $h_{11/2}$) two-neutron quasi-particle (2qp) excitation, while the 7$^-$ states must be built from the ($d_{3/2}$, $h_{11/2}$), ($d_{5/2}$, $h_{11/2}$) and ($g_{7/2}$, $h_{11/2}$) 2qp-excitations.

In order to study on a theoretical level the nuclear renormalization and fragmentation effects, excitations of the next order of complexity must be considered. We have performed selfconsistent shell model calculations [8] in a space that contains, apart from the neutron 2qp-excitations, proton 1p–1h excitations across the Z = 50 gap, including hole shells down to the $f_{7/2}$ shell and coupled excitations of the type (1p–1h) ∗ (2qp). In order to treat correctly intermediate isospin couplings of at least the highest spin states, neutron

Table 1
Transition amplitudes for the lowest 7^-, 8^- and 9^- states in ^{116}Sn. The phase conventions are those appropriate for the program DW81 [12], in which radial wave functions of individual particle states are positive towards infinity. For use in the program WSAXE [14], configurations with a $d_{3/2}$ or $d_{5/2}$ should be used with opposite sign.

p	h	p/n	7_1^-	7_2^-	7_3^-	8_1^-	9_1^-
$h_{11/2}$	$g_{7/2}$	n	0.132	−0.734	−0.319	0.741	−0.795
$g_{7/2}$	$h_{11/2}$	n	0.021	−0.118	−0.051	−0.110	−0.128
$h_{11/2}$	$d_{5/2}$	n	−0.034	0.335	−0.791	−0.350	
$d_{5/2}$	$h_{11/2}$	n	0.004	−0.036	0.086	−0.038	
$h_{11/2}$	$d_{3/2}$	n	−0.442	−0.078	−0.011		
$d_{3/2}$	$h_{11/2}$	n	−0.282	−0.050	−0.007		
$h_{11/2}$	$g_{9/2}$	n	0.008	0.006	0.012	0.019	0.048
$h_{11/2}$	$g_{9/2}$	p	0.024	−0.102	−0.019	−0.035	−0.043
$g_{7/2}$	$f_{7/2}$	p	−0.023	0.014	−0.014		

holes in the $g_{9/2}$-shell were taken into account. As an effective interaction the extended Skyrme force of ref. [9] has been used. This interaction has been shown to describe well the properties of both ground states and excited states throughout the whole mass range with the same parameters. The resulting transition amplitudes for the lowest three 7^- states and the lowest 8^- and 9^- states are given in table 1. The amplitudes

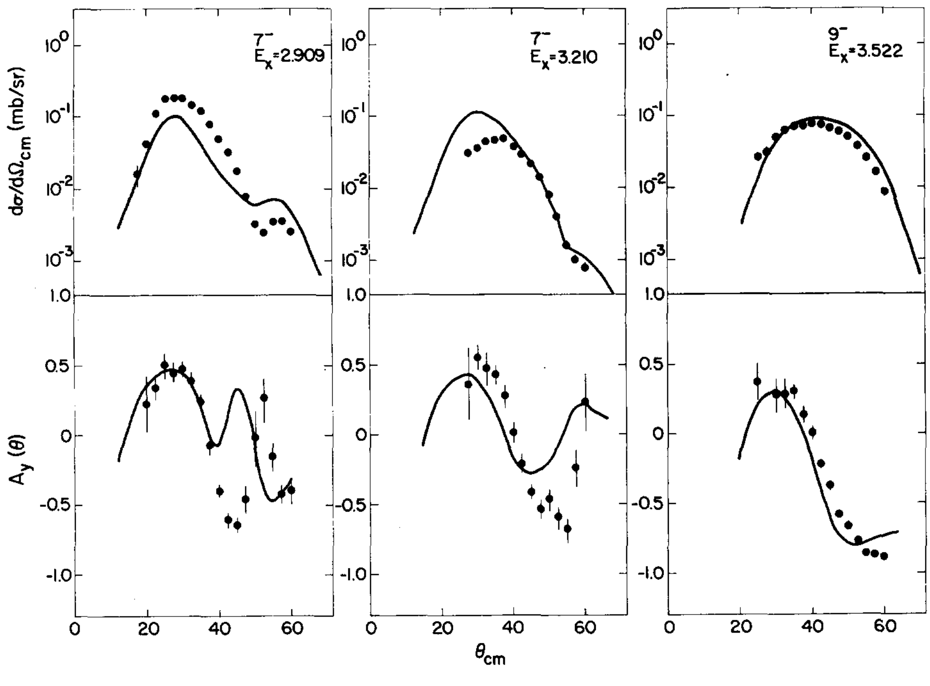

Fig. 2. Cross sections and asymmetries of the ^{116}Sn($\vec{p}$, p$'$)^{116}Sn reaction leading to the $I^\pi = 7^-$ states at E_X = 2.909 and 3.210 MeV and the $I^\pi = 9^-$ state at E_X = 3.522 MeV. The curves have been calculated in DWBA, using the Love–Franey force [11] and the wave functions from the present work. Radial wave functions were generated from a Woods–Saxon potential with r_0 = 1.30 fm and a = 0.65 fm. Depths were adjusted, so as to reproduce the single-particle binding energies.

and relative signs of the neutron configurations agree closely with those derived by Bonsignori et al. [10] in a broken-pair-model calculation, which in contrast to our calculation, does not contain proton excitations and neutron excitations from the $g_{9/2}$ shell.

In fig. 2 the unpolarized cross sections and analyzing powers from the $(\vec{p}, p')$ reaction are shown for the $I^{\pi} = 9^-$ state at $E_x = 3.522$ MeV and the two $I^{\pi} = 7^-$ states at $E_x = 2.909$ and 3.210 MeV. DWIA calculations with the program DW81 [12] ‡1, using the effective N–N t-matrix of Love and Franey [11], are shown that are based on our wave functions. Curves, based on the wave functions of Bonsignori et al. nearly coincide with those in the figure and have therefore not been drawn separately. The states have been identified with the 9_1^-, 7_1^- and 7_2^- model states, respectively. The 7_3^- and 8_1^- states are calculated to have maximum cross sections of only 0.02 mb/sr, which makes it not surprising, that they could not be identified.

For the 9^- state both the cross section and the asymmetry are well reproduced, aside from a slight shift in angle. Both our model calculation and the one of Bonsignori et al. find that the 9^- state is essentially a pure $(g_{7/2}, h_{11/2})$ 2qp-excitation. Adopting this view one finds that a description with the Love–Franey force needs no renormalization.

The cross section of the 7^- state at $E_x = 2.909$ MeV state peaks at a smaller angle than that of the $E_x = 3.210$ MeV state and its asymmetry is more compressed. Although these effects are not exactly reproduced, the fits do reproduce correctly the trends and confirm that the former state is predominantly a $(d_{3/2}, h_{11/2})$ 2qp-excitation, while the latter is mostly of the type $(g_{7/2}, h_{11/2})$. The cross section of the $E_x = 2.909$ MeV is underestimated by the calculations, while that of the $E_x = 3.210$ MeV state is overestimated. However, the sum of their experimental cross sections is very close to the incoherent sum of the calculated cross sections for the $(d_{3/2}, h_{11/2})$, $(d_{5/2}, h_{11/2})$ and $(g_{7/2}, h_{11/2})$ 2qp configurations and arbitrary orthogonal combinations of these configurations preserve this sum. One is thus led to believe that the three 2qp configurations may be fixed slightly differently than in the calculations of refs. [5,7], but that they are otherwise little fragmented.

Using other effective N–N forces one finds these conclusions confirmed, though in general a renormali-

‡1 This program is an extended version of the program DBWA70 [13].

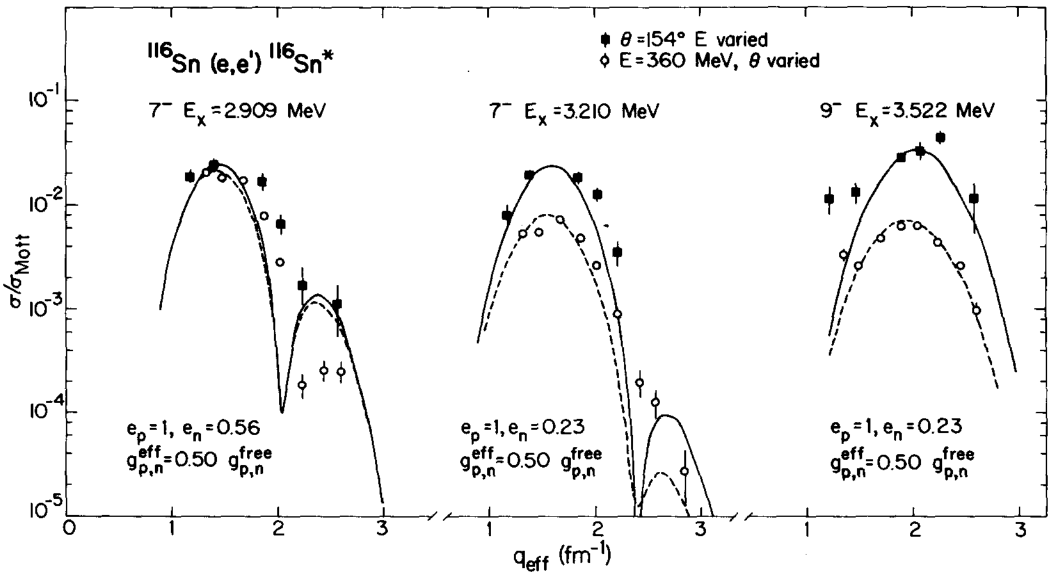

Fig. 3. Experimental form factors from the (e, e′) reaction, compared with curves, based on the model wave functions of this work. The solid curve applies to backward angles ($\theta = 154°$) and the dashed curve to forward angles at $E = 360$ MeV. In the latter case the data cover the angular range 41.5°–96.3°.

zation is needed. These calculations will be described in a forthcoming paper.

In fig. 3 the form factors, obtained from the electron scattering experiment for the same 7^- and 9^- states, are shown. Fits calculated with the programs WSAXE [14] and HEITRA [15], and based on the 7_1^-, 7_2^- and 9_1^- model wave functions are shown. In all three cases effective g-factors $g_{p,n}^{eff}/g_{p,n}^{free} = 0.5$ were used. These values were dictated by the backward angle data points. The effective proton charge was kept equal to its free value, but this choice is not critical, since the contributions from the proton 1p–1h configurations do not dominate the cross sections. For both the 9^- state at $E_x = 3.522$ MeV and the 7^- state at $E_x = 3.210$ MeV $e_n^{eff} = 0.23$ was used. The equality of the value $e_n^{eff} = 0.23$ found from fitting the form factors of the 9^- and 7_2^- states may be incidental. The neutron effective charges, arising from core polarization, or equivalently within the shell-model necessitated by the truncation of the basis space, may be different for each J^π-value, but should, in principle, be equal for states with the same J^π. Fitting the forward angle data of the 7^- state at $E_x = 2.909$ required a larger value of $e_n^{eff} = 0.56$. This may again be taken as an indication that the 7^- states have a slightly different mixing of the 2qp configurations than calculated. The collectivity of the lowest 7^- state is underestimated, as was also indicated by the (p, p′) data.

The proton 1p–1h configurations with the highest spin, expected between 4 and 7 MeV excitation energy, are of the families $(g_{9/2}^{-1}, h_{11/2})$ with $I_{max}^\pi = 10^-$ and $(g_{9/2}^{-1}, g_{7/2})$ with $I_{max}^\pi = 8^+$. Adopting a predictive power for calculations, using the Love–Franey force, one expects for the 10^- state of the first family a cross section of 0.57 mb/sr for the (p, p′) reaction at 45 degrees, if it would exist unfragmented. This would be visible up in the spectrum of fig. 1 as a peak, eight times larger than that of the 9^- state at $E_x = 3.522$ MeV. The unfragmented 8^+ state of the second family is predicted to have a cross section of 0.067 mb/sr, about equal to the 9^- state. Similarly, using $e_p = 1$ and $g_p^{eff} = 0.5\, g_p^{free}$, one finds that such an unfragmented 10^- state should appear in the electron scattering spectrum at $\theta = 154°$ with a peak cross section five times larger than the 3.522 MeV 9^- state, while in the (e, e′) spectrum of fig. 1 the unfragmented 8^+ would be fifteen times larger than the 9^- state. In these (p, p′) and (e, e′) spectra no peak, higher in excitation energy than the 9^- state, comes even close in cross section. Therefore, even the highest spin members of the $(g_{9/2}^{-1}, h_{11/2})$ and $(g_{9/2}^{-1}, g_{7/2})$ families seem to be strongly fragmented.

A complete analysis of the (^{3}He, d) experiment in which data were taken from $\theta = 0°$ to 30° in 2° steps involves fitting some fifty peaks and will be given in a later publication. We emphasize here the salient features of the spectroscopic strength distributions associated with $l_p = 4(g_{7/2})$ and $l_p = 5(h_{11/2})$. This can be illustrated with the difference-spectrum technique. The angular distributions for these l_p-values peak around $\theta = 16°$, where $l_p = 2(d_{3/2}, d_{5/2})$ and $l_p = 0(s_{1/2})$ have pronounced minima. In fig. 4 a difference spectrum of 16° and 6° spectra is shown, in a ratio which leaves only the $l_p = 4$ and 5 strength. In this spectrum nine states or groups of states stand out over

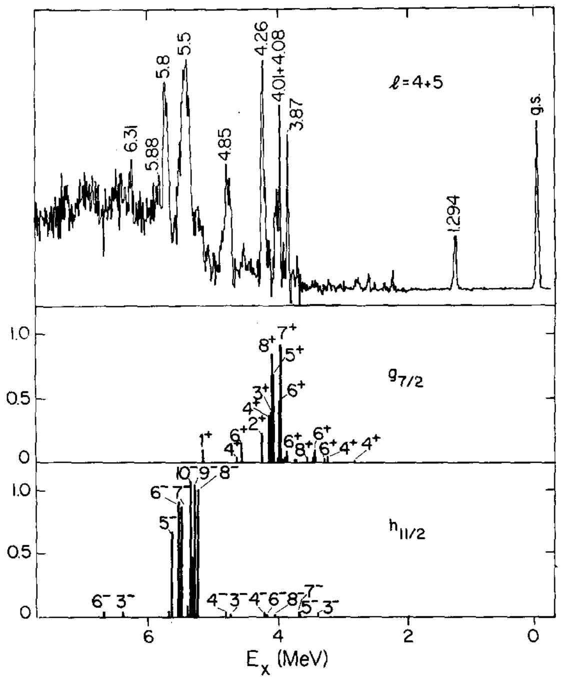

Fig. 4. Difference spectrum (top) of experimental ^{115}In(^{3}He, d)^{116}Sn spectra, taken at $\theta = 16°$ and $\theta = 6°$. The ratio has been chosen to make the $l_p = 2$ strength ($d_{3/2}$ and $d_{5/2}$) vanish, as determined from some rather pure $l = 2$ transitions just below the gap ($E_x = 3.8$ MeV), while also the $l_p = 0$ ($s_{1/2}$) contribution is small. The resulting spectrum thus gives an image of the $l_p = 4$ ($g_{7/2}$) and $l_p = 5$ ($h_{11/2}$) strength. For these latter subshells, strength distributions, resulting from our model calculation [8], are shown (middle and bottom, respectively).

Table 2
States with $E_X > 3.8$ MeV in ^{116}Sn, excited via (α, t) and (^{3}He, d) with l_p = 4 or 5.

E_X (MeV)	l_p	$G(j) = \frac{(2I_f+1)}{(2I_i+1)} C^2S$
3.87	4 (+0)	0.13 (+0.12)
4.00 +4.06	4	1.28
4.26	4	1.32
4.85	4	1.00
5.5	5	1.62
5.8	5	0.84
5.88	5	0.15
6.31	5	0.40

a rather smooth continuum. They are listed in table 2 with their spectroscopic strengths deduced from the full angular distributions. Among them are the states at E_X = 4.85, 5.88 and 6.31 MeV, which are indicated in fig. 1. At high momentum transfer these are the most prominent states in the (p, p′) and (e, e′) spectra, above E_X = 4 MeV. The assignments of the lower states at E_X = 3.87, 4.00, 4.06, 4.26 and 4.85 MeV as $g_{7/2}$ and the groups at E_X = 5.5, 5.8, 5.88 and 6.31 MeV as $h_{11/2}$ have been made on the basis of a comparison with the (α, t) reaction, which, because of its angular momentum mismatch, enhances l_p = 5 over l_p = 4 as compared with the (^{3}He, d) reaction.

The partial sum rule for a definite final spin I_f is $G(j) = (2I_f + 1)/10$ for each of the empty orbitals above the gap. Thus, for instance, unfragmented ($g_{9/2}^{-1}$, $g_{7/2}$; 8^+) and ($g_{9/2}^{-1}$, $h_{11/2}$; 10^-) configurations would have spectroscopic strengths 1.7 and 2.1, respectively. Judging from the transfer data, the best candidate for an 8^+ state with a predominant ($g_{9/2}^{-1}$, $g_{7/2}$) signature would be the state at E_X = 4.26 MeV. However, its cross sections in (p, p′) and (e, e′) are very small, which seems to contradict this assignment.

The strongest clusters of l_p = 5 strength are found at E_X = 5.5 and 5.8 MeV. They are no single peaks, but instead clusters of experimentally unresolved states. In the ($\vec{p}$, p′) and (e, e′) reactions they are hardly distinguished in the spectra, whilst an unmixed ($g_{9/2}^{-1}$, $h_{11/2}$) state, for example, would have a large cross section.

The three states at E_X = 4.85, 5.88 and 6.31 MeV remain as good candidates for high-spin states, but their wave functions cannot be just a single 1p–1h configuration, because they by far do not exhaust a partial sum rule for high spin in the transfer and do not have the expected cross sections in the (p, p′) and (e, e′) reactions.

In conclusion, we have found several high-spin states of 1p–1h signature in ^{116}Sn. The most prominent example is the 9^- state at E_X = 3.522 MeV, which appears to be an almost pure neutron ($g_{7/2}$, $h_{11/2}$) 2qp-excitation. For the $I^\pi = 7^-$ states the neutron 2qp-configurations appear to be mixed but little fragmented. The best candidates for proton 1p–1h high-spin states ad judged on the basis of the stripping reactions, are bad candidates, if judged by their ($\vec{p}$, p′) and (e, e′) cross sections. Conversely, the best candidates from these latter reactions, have too small spectroscopic strength to be pure configurations.

The fragmentation of the l_p = 4 and 5 strength is large: only about 45% of the $g_{7/2}$– and about 25% of the $h_{11/2}$ strength reside in peaks. The rest is spread over very many states and appears as the continuum, over which the peaks stand out. For these high-lying particle states the degree of fragmentation is intermediate between the typical line fragmentation of the valence shells near the Fermi surface and the quasi-continuous strength distribution of the next-higher $h_{9/2}$ and $i_{13/2}$ shells observed in recent proton stripping reactions on ^{144}Sm and ^{116}Sn [16,17]. This situation is reminiscent of the fragmentation of deep hole states. Indeed, in the limit that the $h_{11/2}$ shell is considered empty, ($g_{9/2}^{-1}$, $h_{11/2}$) excitations must be either isoscalar or isovector in character, so that their neutron component is as strong as their proton component. The same states could thus be populated by $g_{9/2}$ neutron pickup on a target nucleus with a single $h_{11/2}$ neutron. Unfortunately, the $11/2^-$ state is never the ground state in stable odd Sn isotopes. However, this argument makes is clear that the spreading widths of the high-lying proton excitations are brought about by the same mechanisms that are responsible for the spreading of deep-lying hole strength.

In fig. 4 the empirical l = 4 + 5 strength is compared with the theoretical $g_{7/2}$ and $h_{11/2}$ distributions obtained from our shell model calculation. In the model the fragmentation of 1p–1h strength is introduced by including the coupling to the 2qp-neutron excitations in the ^{116}Sn basis space. The basis space for ^{115}In accordingly contained coupled configurations of the type

1h ∗ 2qp. It is evident from the figure that the observed fragmentation is much larger than calculated. Since we believe that otherwise our calculation is realistic, we interpret this as evidence that configurations of the next order of complexity: 4qp-excitations, 2p–2h proton excitations and their couplings, are essential in the further spreading of the spectroscopic strength. The actual inclusion of these next-order complex configurations would bring the basis space to prohibitively large dimensions and the strength distribution resulting from our calculation is to be considered as a doorway picture.

The authors wish to thank K. Allaart for many enlightening discussions. This work was supported in part by the "Stichting Fundamenteel Onderzoek der Materie" (FOM) with financial support from the "Nederlandse Organisatie voor Zuiver Wetenschappelijk Onderzoek" (ZWO) and by the US National Science Foundation.

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